Reprinted from Nature, Vol. 316, No. 6027, pp. 418-420, 1 August 1985

Anatomy of a cosmic-ray neutrino source and the Cygnus X-3 system

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There is strong evidence that a compact object in the Cygnus X-3 binary system produces an intense beam of ultra-high-energy cosmic rays. Here, we examine the effects of such a beam hitting the companion star and of the subsequent production of secondary neutrinos. We consider how high a beam luminosity is allowed and how high a neutrino to γ-ray (ν/γ) ratio can be obtained from such a system. We find a maximum allowable beam luminosity of $\sim 10^{42}$ erg s^{-1} for a system consisting of a compact object and a $\sim$1–10 $M_\odot$ main-sequence target star. The proton beam must heat a relatively small area of the target star to satisfy observational constraints on the resulting stellar wind. With such a model, a ν/γ flux ratio of $\sim 10^3$ can result from a combination of γ-ray absorption and a large ν/γ duty cycle ratio. We find that the high density of the atmosphere resulting from compression by the beam leads to pion cascading and a neutrino spectrum peaking at 1–10 GeV energies, which may avoid catastrophic heating of the target star through internal ν interactions. The ν flux and duty cycle are predicted to be accordingly reduced in the energy range above 1 TeV available to a deep underwater neutrino detector.

There has recently been much interest in constructing theoretical models for the Cyg X-3 binary system[1-3] in view of the discovery of ultra-high-energy γ rays from this source[4-6]. The γ-ray flux implies that Cyg X-3 is the first identified source of ultra-high-energy cosmic rays and that the source power is at least 10^{39} erg s^{-1} in cosmic-ray primaries, enough to provide a significant fraction of the 10^{17}-eV cosmic rays in the Galaxy[7]. It has been suggested that Cyg X-3 may be a source of cosmic-ray neutrinos[8,9], produced from the pion decay processes[10] which probably produce the γ rays. This possibility has taken on more interest with the detection of muon signals in deep underground proton decay detectors (ref. 11 and J. Learned, personal communication) from the direction of Cyg X-3 with the 4.8-h periodicity of the Cyg X-3 system. (Such detectors make use of the Cerenkov light emitted by charged particles crossing large volumes of water[12].) The very detectability of such signals, if induced by cosmic-ray neutrinos from Cyg X-3, implies a surprisingly large neutrino flux, primary cosmic-ray beam power ($\gg 10^{39}$ erg s^{-1}) and neutrino-to-γ-ray flux ratio. In the event that the recently detected muons are produced by some other 'X' particle, a large beam power and accompanying neutrino flux are almost certainly implied. Therefore, regardless of the outcome of the present observational situation, it is of interest to consider the theoretical implications of significantly larger beam power.

Hillas[2] has recently presented calculations of γ-ray production from an energetic proton beam, based on a model originally proposed by Vestrand and Eichler[13]. In this model, a beam of $\sim 10^{17}$ eV protons, accelerated at a compact object, impinges on the atmosphere of a companion star, initiating a cascade producing γ rays below 10^{16} eV. The total luminosity of the compact source required to account for the observed γ rays above 10^{12} eV is $L_T = 6.4\times10^{39}\,(\varepsilon_\gamma/0.1)^{-1}(\langle\Omega_B\rangle/4\pi)(0.025/\Delta_\gamma)$ erg s^{-1}, where $\langle\Omega_B\rangle$ is the time-averaged solid angle of the proton beam, ε_γ is the fraction of proton energy converted into γ rays and Δ_γ is the duty cycle of the pulse $\sim$0.025 (ref. 5). Owing to observational uncertainty, a smaller Δ_γ cannot be ruled out, which would imply an even larger beam luminosity. If the part of the beam luminosity striking the target star, L_B, is greater than its Eddington limit,

$$L_{\rm edd} = \frac{4\pi GcM_*}{\kappa} = 1.3\times10^{38}\,\tilde{M}\ \text{erg s}^{-1} \tag{1}$$

where M_* is the target star mass and κ is the electron scattering opacity, a mass loss will result. (Note that all quantities with overbars will be in solar units, for example, $\tilde{R} = R/R_\odot$). The mass loss from the system is limited by the observed derivative, $\dot{P}$, of the 4.8-h orbital period[14], through the relation[15]

$$\dot{P}/P = 2\dot{M}/M_T \tag{2}$$

where M_T is the total mass of the system and $\dot{M}$ is the mass loss rate from the target star (assuming that the orbital angular momentum is lost from the system and that the orbit decay from gravitational radiation is small). When the beam hits the atmosphere of the target star, roughly half of the beam energy flux will be thermalized in the stellar atmosphere, producing an outward radiative flux. The resulting radiation pressure can drive a wind with a terminal velocity v given by

$$v^2 = \frac{2}{R_*}\left(\frac{L_B\kappa}{4\pi fc} - GM_*\right) \tag{3}$$

where R_* is the target star radius and f is the fraction of the area of the star illuminated by the proton beam. Energy conservation requires that

$$L_B \simeq L_{\rm rad} + \frac{GM_*\dot{M}}{R_*} + \tfrac{3}{2}\dot{M}\left(\frac{kT}{m}\right) + \tfrac{1}{2}\dot{M}v^2 \tag{4}$$

where m is the mean atomic mass of the atmosphere. The radiative flux, $L_{\rm rad}$, can be neglected if $L_B \gg L_{\rm edd}$, as can the thermal energy flux (third term). The mass loss rate, combining equations (3) and (4), is then

$$\dot{M} = 4\pi fR_*c/\kappa = 9.6\times10^{-4}f\tilde{R}\,M_\odot\,\text{yr}^{-1} \tag{5}$$

The mass loss rate is independent of the beam luminosity and the mass of the target star (provided that $L_B > L_{\rm edd}$). Using equations (2) and (5), we can derive a relation between R_* and M_T: $R_* = 1.5\times10^8 f^{-1}\tilde{M}_T$ cm. This relation is valid only in the case where the wind from the target star is lost from the system. In this case, the escape velocity of the system is less than v and $f < 12\tilde{R}^{-1}\tilde{M}_T^{-3/2}L_{39}$. Therefore, since $f<1$ and $L_B > 10^{39}$ erg s^{-1}, nearly all of the wind escapes the system. Equations (2) and (5) may then be used to find the fraction of the area of the star heated by the beam such that the resulting mass loss gives the observed $\dot{P} = 1.2\times10^{-9}$ (ref. 14):

$$f = 2.2\times10^{-3}\tilde{M}_T\tilde{R}^{-1} \tag{6}$$

The angular extent θ of the time-averaged beam in one direction is given by $\tan\theta = 2\pi f/(a/R_* - 1)$. Since a/R_* may be very close to unity (as indicated from the X-ray light curve), the small area predicted by equation (6) does not necessarily imply a small beam angle. The X-ray light curve shows deep but not sharp eclipses[16,17], thus the optical depth of the outflowing material to electron scattering is $\tau_x \sim 1$ at distances of the order of the orbital separation, $a = 10^{11}\tilde{M}_T^{1/3}$ cm. The optical depth τ_x at a distance r is $\tau_x = \kappa\rho r$, where ρ is the mass density of the wind, assuming that v is constant for $r \gg R_*$. Then using equation (6), and mass conservation in the wind, $\dot{M} = 4\pi fr^2\rho v$, the radius at which $\tau_x = 1$ is $R_1 = (c/v)R_*$ and

$$\frac{R_1}{a} = 5.3\tilde{M}_T^{1/6}\tilde{R}L_{39}^{-1/2} \tag{7}$$

Thus, the condition $R_1/a \sim 1$ required by the X-ray light curve and suggested by cocoon models for Cyg X-3 (refs 16, 17) can be satisfied.

When the beam grazes the stellar limb, it can temporarily heat an area of $\sim 4\pi fR_*^2$ to a temperature of

$$T = 6\times10^5\tilde{M}_T^{-1/4}\tilde{R}^{-3/4}L_{39}^{1/4}\ \ \text{K} \tag{8}$$

This increases the atmospheric scale height to

$$h = kT/mg = 2\times10^9 L_{39}^{1/4}\tilde{M}_T^{-1/4}\tilde{R}^{5/4}\tilde{M}^{-1}\ \ \text{cm} \tag{9}$$

where g is the surface acceleration due to gravity. This is a lower limit on h, as radiation pressure will tend to increase the scale height. The primary protons have a mean free path for interaction of $\Lambda_p \sim 30\ \mathrm{g\ cm^{-2}}$ at high energies[18]. They produce π^0 mesons which decay into γ rays having a mean free path against pair production off electrons in the stellar atmosphere of $\sim 50\ \mathrm{g\ cm^{-2}}$ at an energy of ~ 1 TeV. Because of absorption effects, the γ rays are strongly produced only at the limb of an atmosphere[19]. The primaries also produce charged pions which decay into neutrinos. The efficiency for neutrino production reaches a plateau[9] at an atmospheric depth of $5\Lambda_p$.

Whereas in the case where the beam is at the stellar limb, heating of the atmosphere increases the scale height h, in the case where the beam hits the star directly, ram pressure from the beam with energy flux F_B, acting downward in concert with gravity, decreases h. In this case, h is given by

$$h = \frac{kT}{m}\left[\frac{F_B}{5\Lambda_p c}\delta + \frac{GM_*}{R_*^2}\right]^{-1} \simeq 3\times 10^7 L_{39}^{-3/4}\tilde{R}^{1/4}\tilde{M}_T^{3/4} \quad \mathrm{cm} \tag{10}$$

where we have assumed energy and momentum balance in the atmosphere. The parameter $\delta = (1 - \Phi_u/\Phi_T)$, where Φ_T is the total radiative flux and Φ_u is the upward flux component, takes values between 0 and 1, depending on geometry and radiation transfer in the atmosphere. When $\delta = 0$, all of the radiation propagates upwards, producing a force which directly opposes the ram force, and no compression occurs. When $\delta = 1$, all radiation propagates or escapes sideways and the atmosphere will be compressed by the full ram force of the beam. However, at the edges of the beam, the escaping radiation should drive the mass loss estimated by equation (5). Since the smallest dimension of the beam, $x = 2\pi f R_* = 9\times 10^8$ cm, is less than the uncompressed atmospheric scale height as derived in equation (9), we can assume that $\delta \sim 1$. Because the beam encounters much higher grammage here, the γ rays are completely absorbed. For ν energies below ~ 100 TeV, the cross-section for interactions with nucleons is in the linear region[10] and is given by (expressing all energies in TeV units) $\sigma_\nu = 7\times 10^{-36} E_\nu\ \mathrm{cm^2}$. The optical depth of a star to neutrinos is then given by $\tau_\nu = 0.41 E_\nu \tilde{M}\tilde{R}^{-2}$, so that the critical energy above which the star is optically thick to neutrinos is $E_\nu^* = 2.4\tilde{M}^{-1}\tilde{R}^2$TeV.

If neutrinos are generated above E_ν^*, such neutrinos will be absorbed within the main body of the star uniformly, converting into high-energy muons and electrons which heat up the inside of the star where the photon diffusion escape time is $\sim 10^6$ yr. Absorption of an energy flux significantly above the Eddington limit would result in disruption of the target star on a timescale $<10^3$ yr. Therefore, such neutrinos must not be generated in a plausible theoretical scenario.

Fortunately, ν production above ~ 1 TeV energy can be suppressed. Owing to relativistic time dilation, the distance travelled by the parent pion before decay as a function of energy is given by $l_\pi = 5.57\times 10^6 E_\pi$ cm. If this distance is greater than the interaction mean free path, the pion will interact and produce a pion cascade rather than decay into neutrinos directly[20]. This will result in the suppression of higher-energy neutrinos and the multiple production of neutrinos of low enough energy (from the decay of lower-energy parent pions) that they can travel through the star without being absorbed. This situation will occur if the atmospheric density at a depth of ~ 50–$200\ \mathrm{g\ cm^{-2}}$ (roughly where ν production is most efficient) is greater than $(l_\pi \sigma_\pi)^{-1}$ where σ_π (≥ 1 TeV) $\simeq 6\times 10^{-26}\ \mathrm{cm^2}$ (ref. 18), corresponding to a critical energy, $E_\nu^c \simeq (10^{-6}\ \mathrm{g\ cm^{-3}}/\rho)$ TeV. For $E_\nu^c < E_\nu^*$, we find the required density criterion for the neutrino production region should be $\rho < \rho_c = 4.2\times 10^{-7}\ \mathrm{g\ cm^{-3}}\ \tilde{M}\tilde{R}^{-2}$. From the scale height in equation (10), we have $\rho \simeq 5\Lambda_p/h = 5\times 10^{-6} L_{39}^{3/4}\tilde{M}_T^{-3/4}$. This effectively suppresses the more dangerous νs since $E_\nu^c < E_\nu^*$. Because the energy threshold for ν detection with DUMAND[21] is ~ 1 TeV, we predict that Cyg X-3 may not be a superstrong source for such a detector.

The fraction of π-decay νs that interact even when $\tau_\nu \ll 1$ will contribute to heating the stellar interior. The amount of energy absorbed by the star is $\varepsilon_\nu \tau_\nu L_B \sim 4.1\times 10^{38} \varepsilon_\nu L_{39}\tilde{M}\tilde{R}^{-2}\langle E_\nu\rangle\ \mathrm{erg\ s^{-1}}$, where ε_ν is the efficiency for neutrino production. The average energy $\langle E_\nu\rangle \simeq 1$–$10$ GeV for $\rho = 10^{-6}\ \mathrm{g\ cm^{-3}}$ (ref. 20). In this case, detailed cascade calculations required to obtain a reliable heating estimate may impose severe constraints on the Hillas model.

There is a limit to high-energy ν suppression. There will always be high-energy νs from prompt decay of mesons carrying charm or higher mass quark flavours. The ratio ξ of forward production of these mesons to πs is probably $\sim 10^{-4}$–10^{-3} at high energies, although the experimental situation is still in some doubt[22]. Taking $\xi \geq 10^{-4}$, we can use the existence of prompt neutrinos to place a theoretical upper limit on the beam power, L_B.

Let us assume that the π-decay neutrinos are effectively suppressed at high energies by ram-pressure density enhancement in the region of the stellar atmosphere directly below the particle beam where the πs are produced (see equation (10)). We are then still left with the prompt νs which will give a lower limit to the high-energy ν flux. As about half the beam energy goes into first generation νs from πs (or their cascade equivalent), taking $\xi \geq 10^{-4}$ implies that at least $\sim 5\times 10^{-5} L_B$ ($>E_p^*$) will go into prompt νs that heat the stellar interior, where $E_p^* \simeq 4E_\nu^*$ and E_ν^* is the energy above which σ_ν is large enough to assure $\tau \geq 1$. For main-sequence stars of $1M_\odot$ and $10M_\odot$, $E_p^* \sim 10$ and ~ 25 TeV respectively. In the Hillas[2] model, essentially all of the beam power lies in this energy range, so that the interior heating rate from prompt νs is given by $\eta_{pr} \geq 5\times 10^{-5} L_B$. If $\eta_{pr} > L_{edd} \simeq 10^{38}\ \tilde{M}\ \mathrm{erg\ s^{-1}}$, the target star would be disrupted. This places an upper limit on the beam power for the Cyg X-3 system:

$$L_B \leq 2\times 10^{42}\tilde{M} \quad \mathrm{erg\ s^{-1}} \tag{11}$$

Unless $\eta_{pr} < L_*$, the intrinsic stellar luminosity, the structure of the star will be modified and its radius will expand. For main-sequence stars, this condition is given by

$$L_B \leq \begin{cases} 8\times 10^{37} & \mathrm{erg\ s^{-1}}, \quad \tilde{M} = 1 \\ 2\times 10^{40} & \mathrm{erg\ s^{-1}}, \quad \tilde{M} = 4 \\ 4\times 10^{41} & \mathrm{erg\ s^{-1}}, \quad \tilde{M} = 10 \end{cases} \tag{12}$$

This can be compared with the π-decay neutrino heating limit, $\eta_\pi \equiv \varepsilon_\nu \tau_\nu L_B < L_*$,

$$L_B \leq \begin{cases} 10^{38}(\varepsilon_\nu/0.1)^{-1}\langle E_\nu/10^{-3}\rangle^{-1} & \mathrm{erg\ s^{-1}}, \quad \tilde{M} = 1 \\ 4.7\times 10^{40}(\varepsilon_\nu/0.1)^{-1}\langle E_\nu/10^{-3}\rangle^{-1} & \mathrm{erg\ s^{-1}}, \quad \tilde{M} = 4 \\ 1.3\times 10^{42}(\varepsilon_\nu/0.1)^{-1}\langle E_\nu/10^{-3}\rangle^{-1} & \mathrm{erg\ s^{-1}}, \quad \tilde{M} = 10 \end{cases} \tag{13}$$

Should either prompt neutrino or π-decay heating exceed the stellar luminosity, radiative transfer will no longer be efficient enough to carry the increased luminosity to the surface and the star will convectively transport this additional energy. The star will expand (see ref. 23), returning to the Hayashi track in the Hertzsprung–Russell diagram, until cooling to a surface temperature of ~ 3 to 4×10^3 K. The new radius which results is $R_* \sim 7\times 10^{13}\eta_{39}^{1/2}\ (T/4{,}000\ \mathrm{K})^{-2}$ cm. This R_* would then exceed the stellar orbit, and may quench the beam of the compact object, reducing the radius until it is within the size of the orbit. The stability of such a system is unknown.

We now consider the ratio of ν-to-γ-ray fluxes in the case where the above constraints on L_B are satisfied. There are two situations:

(1) In the case where the beam is grazing the target star's atmosphere at the limb, for pathlengths $X \leq 25\ \mathrm{g\ cm^{-2}}$, comparable fluxes of ν and γ rays are produced[10]. For astrophysically significant production pathlengths $\sim 60\ \mathrm{g\ cm^{-2}}$, the ν/γ ratio is considerably higher owing to attenuation of the γ rays by pair production. This flux ratio is approximately an order of magnitude for a cosmic-ray primary spectrum of index near 2 (ref. 9). The fraction of the orbital period during which the stellar limb eclipses the source of the beam will be denoted as Δ_a.

(2) Where the beam is hitting the star directly, the neutrinos

produced by the decay of the pions within $X = 5\Lambda_p$, pass through the star and exit provided they are below some critical energy E^*_ν. Denoting the fraction of the period in which the beam is directly hitting the star by Δ_*, the ratio of neutrino flux to γ-ray flux per cycle is given by $F_\nu/F_\gamma = 10\Delta_*/\Delta_a$, where the factor of 10 is due to the ratio of ν-to-γ-ray production efficiencies. The ratio, Δ_*/Δ_a is given by $\Delta_*/\Delta_a \simeq R_*/h \simeq 100$, with $R_* = \sim 10^{11}$ cm and $H = 10^9$ cm (see equation (9)). Therefore, the ratio of neutrinos to γ rays per orbit can be as high as $\sim 10^3$. Thus, one can obtain much higher ν/γ ratios than previously thought.

Our results therefore indicate that high ν/γ ratios and neutrino fluxes are possible from a source such as Cyg X-3, but that there are also strict theoretical constraints from such models. Such constraints are independent of the details of the models involved and should be useful in considering future observational data.

Note that since submission of this letter two papers[24,25] on neutrino production in Cyg X-3 have come to our attention.

We thank T. Gaisser, F. Halzen, D. Kazanas, J. Learned, M. Marshak and R. Ramaty for helpful discussions. J.J.B. is a National Academy of Sciences National Research Council Postdoctoral Resident Research Associate.

Received 11 April; accepted 12 June 1985.

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